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# A Simple Complete Model of Gauge-Mediated SUSY-Breaking and Dynamical Relaxation Mechanism for Solving the $\mu$ -Problem

S. Dimopoulos<sup>(1),(2)</sup>, G. Dvali<sup>(1)</sup> and R. Rattazzi<sup>(1)</sup><sup>(1)</sup> *Theory Division, CERN**CH-1211 Geneva 23, Switzerland*<sup>(2)</sup> *Physics Department, Stanford University**Stanford, CA 94305, USA*

## Abstract

In this paper two things are done. First, we propose a simple model of dynamical gauge-mediated SUSY breaking. This model incorporates a dynamical relaxation mechanism which solves the  $\mu$ -problem with no light fields beyond those of the MSSM. In the second part of the paper we show how this mechanism is generalized and give two more examples in which it is incorporated.

## A. A Simple Model

Gauge-mediated (GM) theories of supersymmetry breaking [1] have recently attracted a great deal of attention because they provide a natural solution to the supersymmetric flavour problem [2] and have new phenomenological signatures. Since the original pioneering work [3], there has been significant progress in building models which combine dynamical supersymmetry breaking with gauge mediation. (see e.g. [4] - [13]). In particular, new mechanisms based on quantum-modified moduli spaces [15] have been explored [14]. It has also been shown that a class of simple strongly coupled gauge theories, which usually preserve supersymmetry, do break it dynamically when combined with the stringy anomalous  $U(1)$  [16].

We have recently proposed a class of simple theories in which SUSY is unbroken in the global minimum, and yet the soft terms are dynamically generated because we live in a metastable SUSY-breaking plateau [10]<sup>1</sup>. Unfortunately the most elegant of our theories had negative mass<sup>2</sup> scalars [17]. In this paper we propose a simple class of realistic theories with such metastable vacua. These theories provide for a simple solution to the  $\mu$ -problem based on a dynamical relaxation mechanism [18]; furthermore they do not require the introduction of any input mass parameters. In the second part of the paper we show how to generalize this mechanism and give two more examples incorporating it.

According to the general mechanism of ref. [10] our gauge group includes three factors  $SU(5)_S \otimes SU(5)_W \otimes G_B$ . Here  $SU(5)_W$  is the ordinary (weakly coupled) grand unified group which accommodates quarks and leptons in three families of  $(1, 10, 1) + (1, \bar{5}, 1)$  representations. For our discussion it is not essential that  $SU(5)_W$  be realized as a gauge symmetry of the theory above the GUT scale. In particular, the  $SU(3)_C \otimes SU(2) \times U(1)_Y$  subgroup would suffice. Nevertheless, for convenience, we will classify fields by  $SU(5)_W$  representations.  $G_B$  is a weakly coupled ‘balancing group’, whose role is to stabilize the supersymmetry-breaking minimum by a mechanism similar to that of Witten’s ‘inverse

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<sup>1</sup>One of our models has a significant ideological overlap, including the problem of tachyonic sfermions, with ref [9]. An interesting model along similar lines was recently proposed in [11].

hierarchy' scenario [19]. This role can be played by any asymptotically free gauge group with some vector-like matter content. For definiteness we will take  $G_B = SU(2)_B$  with a single flavour  $\psi, \bar{\psi}$ . Finally,  $SU(5)_S$  is a strongly coupled group with five flavors  $Q_i \bar{Q}^i$ , where  $i = 1, \dots, 5$  is a flavour index which corresponds to  $SU(5)_W$ . Thus the 'quarks' of the strongly-coupled group are in the representation  $(5, \bar{5}, 1) + (\bar{5}, 5, 1)$ . We will denote the strong scale of  $SU(5)_S$  by  $\Lambda$ . To complete the matter content, we introduce four additional gauge singlets  $Y, I, S$ , and  $N$ . Their role will become clear later.

Our classical superpotential is

$$W = \alpha Y \text{Tr} M + I(\beta Y^2 - \gamma \psi \bar{\psi}) + S(\lambda_1 H \bar{H} + \lambda_2 N^2/2 - \lambda \text{Tr} M) \\ + \text{standard MSSM (GUT) superpotential} \quad (1)$$

where  $M_i^k = Q_i \bar{Q}^k$  are the 'mesons' of the strongly coupled group and  $H, \bar{H}$  are the MSSM Higgs doublets.

The role of the various terms in the superpotential can be briefly summarized as follows.

- 1) After accounting for the non-perturbative dynamics, the first term produces a supersymmetry breaking plateau along which  $F_Y \sim \Lambda^2$  and  $Y$  is undetermined.
- 2) The second term is the main improvement with respect to the models of ref. [10]. It relates the vacuum expectation value (VEV) of the singlet  $Y$  to the VEV of  $\psi \bar{\psi}$ . Thus the balancing group is broken ('Higgsed') along the plateau and one-loop effects can create a stable minimum at  $Y \gg \sqrt{F_Y}$  via Witten's inverse hierarchy mechanism. In this minimum  $Q, \bar{Q}$  automatically play the role of the usual messengers of supersymmetry breaking. Note that without the second term, the only way to stabilize the SUSY breaking minimum is with a gauge-non-singlet  $Y$  as it was done in ref. [10]. This would make it more difficult to solve the  $\mu$  problem and may also lead to negative sfermion mass<sup>2</sup> [17].
- 3) The last term solves the  $\mu$  problem via the mechanism of ref. [18]: the  $\mu$  and  $B\mu$  terms are induced from the original SUSY-breaking scale  $F_Y$  as one- and two-loop effects respectively.

Let us now discuss this model in more detail. First we study its classical vacuum manifold (the moduli space). It includes a flat direction that can be parameterized by  $Y$ . Along this direction the  $Q, \bar{Q}$  states get masses  $\sim Y$ . For  $\beta = \gamma = 0$  there would be two other independent flat directions on the moduli space: a trivial one, given by the free field  $I$ , and another, parameterized by the invariant  $\phi = \sqrt{2\psi\bar{\psi}}$ . The latter is the only flat direction allowed by the  $SU(2)_B$   $D$ -terms. Along it,  $SU(2)_B$  is completely Higgsed and its gauge superfields get a mass  $= g_B\phi$  by eating up the three complex degrees of freedom from  $\psi, \bar{\psi}$ . Thus, in the  $\beta = \gamma = 0$  limit we would have a moduli space of complex dimension 3. However, for  $\beta, \gamma \neq 0$  two out of three flat directions are lifted. This is because the  $F_I$  and  $F_\phi$  flatness conditions force to have

$$Y\sqrt{\beta} = \phi\sqrt{\gamma/2} \text{ and } I = 0 \quad (2)$$

As a result  $I$  and the combination  $Y\sqrt{\beta} - \phi\sqrt{\gamma/2}$  get mass  $= 2\sqrt{\beta^2 + \beta\gamma/2}\langle Y \rangle$ . Notice that we define  $\beta$  and  $\gamma$  to be positive and real. The other superposition remains flat and we parameterize it by  $X = Y\sqrt{1 + 2\beta/\gamma}$ . Thus, perturbatively the massless degrees of freedom along this branch of the moduli space are: the  $SU(5)_S$  vector supermultiplet, the chiral superfields  $X, S$  and  $N$ , plus the ordinary MSSM spectrum (quarks, leptons etc.). Integrating out the heavy modes, the perturbative effective low energy superpotential is

$$W = S(\lambda_1 H\bar{H} + \lambda_2 N^2/2) + \text{standard MSSM (GUT) superpotential} \quad (3)$$

We see that  $X$  simply dropped out of the superpotential. The Kähler potential for  $X$  is determined by the wave functions of the fields in the original theory as discussed in ref. [10]

$$K(X, S) = \left( yZ_Y(XX^\dagger) + (1 - y)Z_\psi(XX^\dagger) \right) XX^\dagger + \left( Z_{YS}(XX^\dagger)\sqrt{y}SX^\dagger + \text{h.c.} \right) \quad (4)$$

where  $y = \gamma/(\gamma + 2\beta)$ . Notice that the wave function mixing between  $Y$  and  $S$ , which is radiatively induced by the Yukawa couplings, leads to a corresponding term for  $X$ . The crucial point is that, while  $Z_Y$  is only renormalized at 1-loop by matter diagrams,  $Z_\psi$  is also corrected by loops involving the massive  $SU(2)_B$  vector superfields. By defining the Yukawa couplings at a renormalization scale  $\mu_R$  at which  $Z_\psi = Z_Y = 1$ , the first term in brackets in eq. 4 reads at lowest order in  $\ln(X/\mu_R)$

$$K \simeq XX^\dagger \left( 1 - \frac{(25\alpha^2 + 4\beta^2)y + (\gamma^2 - C_B g_B^2)(1-y)}{16\pi^2} \ln \frac{XX^\dagger}{\mu_R^2} \right) \quad (5)$$

where  $C_B$  is the Casimir of  $\psi$ , equalling  $4/3$  for  $SU(2)_B$ . Now let us discuss supersymmetry breaking. It occurs due to the non-perturbative  $SU(5)_S$  dynamics which generates a superpotential linear in  $X$ . The source for this effect is gaugino condensation in  $SU(5)_S$  which for  $X \gg \Lambda$  behaves as pure super-Yang-Mills theory. Gaugino condensation induces the non-perturbative contribution to the superpotential [22,23]

$$W = \langle \bar{\lambda}\lambda \rangle = \Lambda_L^3 \quad (6)$$

where  $\Lambda_L$  is the low energy scale of the  $SU(5)_S$  theory. This scale is  $X$  dependent, because so are the masses of the quarks we have integrated out. Matching the low and the high energy gauge couplings [24] we get the effective superpotential for  $X$

$$W = \alpha Y \Lambda^2 = \rho X \Lambda^2 \quad (7)$$

where  $\rho = \alpha\sqrt{y}$ . This superpotential breaks supersymmetry spontaneously:  $F_X = \rho\Lambda^2$ . After taking into account the Kähler renormalization the effective potential for  $X$  at one loop reads

$$V = K_X^{X-1} W_X W^{X*} = \frac{\rho^2 \Lambda^4}{y Z_Y(XX^\dagger) + (1-y) Z_\psi(XX^\dagger)} \quad (8)$$

The balancing between gauge and Yukawa coupling contributions in  $Z_Y$ ,  $Z_\psi$  can create local minima in  $V$ . This can be seen explicitly by using the leading terms in eq. 5. For a range of parameters, as it was the case in ref. [10], the local minimum is at  $X \gg \Lambda$ , where our perturbative approximation to  $K$  is reliable. In this minimum the messengers have masses  $\sim X$  with fermi-bose mass splitting due to  $F_X \neq 0$ . As a result, all the conditions of standard gauge-mediation are satisfied and supersymmetry breaking is transmitted to the MSSM sector after integrating out  $\bar{Q}, Q$ .

Now let us show how the  $\mu$  and  $B\mu$  terms are generated after supersymmetry breaking. This is arranged by the second term in (1) through the mechanism of [18]. Integrating the heavy mesons out and substituting  $\text{Tr} M = \Lambda^2$  we get the effective superpotential

$$W = S(\lambda_1 H \bar{H} + \lambda_2 N^2/2 - \lambda \Lambda^2). \quad (9)$$

This has precisely the form introduced in ref. [18]. In the limit  $Z_{YS}(\langle X \rangle) \sim 0$ , the field  $X$  is decoupled (see eq. 4) from the above sector in the effective theory below  $\langle X \rangle$ . This situation is precisely the one assumed in ref. [18] and we will do the same in the following for the sake of simplicity. The conclusions are qualitatively unchanged in the general case of  $Z_{YS}(\langle X \rangle) \sim 1$ . The dynamics that leads to nonzero values for  $\mu$  and  $B\mu$  can be described in three steps.

a) The superpotential in eq. 9, induces at tree level the VEV for the  $N$  field  $N^2 = \frac{2\lambda}{\lambda_2} \Lambda^2$ .

The other fields  $(S, H, \bar{H})$  stay zero to this order<sup>2</sup>

b) At one-loop a tadpole term for the scalar component of  $S$  is induced [18]

$$-S \frac{25\lambda\rho^3}{16\pi^2} \Lambda^4 / \langle X \rangle + \text{h.c.} \quad (10)$$

(According to the rules of ref. [17], this term can be seen to arise from the renormalization of  $Z_{YS}$ ). This plays a crucial role, since it generates the VEV  $S = -\frac{25\rho^3}{32\pi^2} \Lambda^2 / \langle X \rangle$  and thus the  $\mu$  term! A  $B\mu$  term is also generated, since for  $S \neq 0$  the cancellation of  $F_S$  is impossible and one gets  $B\mu = \lambda_1 F_S = -\frac{\lambda_1}{\lambda_2} \left( \frac{25\rho^3}{32\pi^2} \right)^2 \Lambda^4 / \langle X \rangle^2$ .

c) At two-loops there is an additional contribution to the soft mass of  $N$ , which shifts the one-loop value of  $B\mu$  so that finally the total expressions for  $\mu$  and  $B\mu$  are [18]

$$\begin{aligned} \mu &= \frac{25\rho^3\lambda_1}{32\pi^2\lambda_2} \Lambda^2 / \langle X \rangle \\ B\mu &= -\frac{50\lambda^2\lambda_1\lambda_2\rho^2}{(16\pi^2)^2} \left( 1 + \frac{25\rho^4}{8\lambda_2^2\lambda^2} \right) \Lambda^4 / \langle X \rangle^2 \end{aligned} \quad (11)$$

Some remarks are now in order. The first concerns the use of a singlet  $Y$  rather than a non-singlet  $X$ , as done in ref. [10], to give a mass to the messengers  $Q\bar{Q}$ . The reason for this choice is that we want the same invariant of the strong group to couple to both

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<sup>2</sup>For  $S = 0$  there is a continuous degeneracy of the vacuum parametrized by the holomorphic invariants  $N^2$  and  $H\bar{H}$  subject to the constraint  $F_S = 0$ . However, since at one-loop  $S \neq 0$  (see the text below), the minimum is fixed at  $H = \bar{H} = 0$  for a range of parameters [18].

$X$  and  $S$ , otherwise a linear term in  $S$  is not generated in  $W_{eff}$ . For instance one may try and use the  $SU(6)^2$  model of ref. [10] and put  $Y$  in the adjoint of  $SU(6)_B$ . In this case however  $\langle M \rangle$  is also in the adjoint, so that  $\text{Tr}\langle M \rangle = 0$ . The use of a singlet  $Y$ , and the addition of the  $I, \psi, \bar{\psi}$  sector, bypasses this difficulty.

The second remark concerns the genericity of our mechanism. The superpotential in (1) is not the most general one could write down. However the results for  $\mu$ ,  $B\mu$  are qualitatively unchanged in a vast class of models of this type. The only crucial feature of eq. (1) which must be preserved is that  $H$ ,  $\bar{H}$  and  $N$  appear only via the expression  $H\bar{H} + N^2$ , as this is what allows  $B\mu$  to be a dynamical quantity. In the model we discussed, this property cannot be enforced by any symmetry. However, as suggested in ref. [18], one can conceive a scenario where  $N^2 \rightarrow N\bar{N}$  so that  $H\bar{H} + N\bar{N}$  is enforced by an  $SU(3)$  global symmetry. Such an effective symmetry of the superpotential can indeed result as a subgroup of an original  $SU(6)$  GUT group in a class of models that naturally solve the doublet triplet splitting problem [20]. With this assumption, eq. (1) can be safely generalized by imposing an (anomalous)  $R$  symmetry under which the charges of  $(Y, S, H\bar{H}, M, \phi^2, I)$  are  $(2, 2, 0, 0, 4, -2)$ . Notice that  $S$  and  $Y$  have the same quantum numbers, but we can always define  $Y$  to be the combination that couples to  $M$  and not to  $H\bar{H}$ . Thus, in the most general renormalizable  $W$  only the intermediate piece is modified to

$$I(\beta_1 Y^2 + \beta_2 S^2 + \beta_3 YS - \gamma \psi \bar{\psi}). \quad (12)$$

This preserves our mechanism: at  $S = 0$  there is still a flat direction  $X$  along which supersymmetry is broken. The only difference with respect to the previous case, is that the field that gets a mass by pairing up with  $I$  is now a linear combination of  $Y$ ,  $\phi$  and  $S$ .

To conclude, notice that we have neglected the non-perturbative effects of the balancing group. This is because the stabilization mechanism depends on a power of  $g_B$  so that it works in the perturbative regime. Non perturbative effects, on the other hand, go like  $\exp - (8\pi^2/g_B^2)$  and become quickly negligible. In our  $SU(2)_B$  example there is an instanton contribution to  $W_{eff} \sim \Lambda_B^5/\phi^2$ . This effect restores supersymmetry somewhere along the  $X$  line. However, (as discussed in ref. [10] in a similar case), it is consistent for

us to neglect this term as long as it does not perturb the local minimum where we live. A direct estimate shows that this is the case already for  $g_B(X) \leq 1$ . Moreover one can imagine models where the balancing group dynamics does not generate a superpotential. An example is  $SU(2)_B$  with matter content just given by a triplet  $T$  with  $\psi\bar{\psi} \rightarrow T^2$  in eq. (1).

### B. Generalizing the Vacuum Relaxation Mechanism for $\mu$

The typical problem in GM theories is that  $\mu$  and  $B\mu$  are induced at the same loop order resulting in

$$B\mu \gg \mu^2. \quad (13)$$

In the previous section we have shown a concrete GM model where a dynamical relaxation mechanism can avoid this problem. In this section we will generalize this mechanism and derive sufficient conditions for it to work. The key point is that both  $\mu$  and  $B\mu$  depend on dynamical fields which, after supersymmetry breaking, will adjust to the right vacuum expectation values due to energy minimization. We note however, that all these additional fields must have masses of order  $F_X$ , so that below this scale the low energy spectrum is just that of the MSSM (plus possibly some additional light states, like  $X$ , which only interact via  $\langle X \rangle^{-1}$  suppressed couplings). This is the crucial difference from the conventional singlet models [3] in which  $\mu$  is induced as a VEV of a *light* singlet field (an example of the later is the ‘sliding singlet’ mechanism for generating the  $\mu$  term in GM theories [21]). The necessary ingredient of our approach is the gauge-singlet field  $S$  that couples to the Higgs doublets in the superpotential (with coupling constant of order one, which we do not display for the time being).

$$W = SH\bar{H} \quad (14)$$

The expectation values of  $S$  and  $F_S$  will then respectively set the scale of  $\mu$  and  $B\mu$ . Following the notations of the previous section, let  $X$  be a superfield that couples to the messengers  $(\bar{Q}, Q)$  and breaks SUSY, through a nonzero  $F_X$  term. Then, to be in the right ball park the, the VEV of  $S$  should be induced as a one-loop effect from the



scale  $\frac{F_X}{\langle X \rangle}$ . For this to be the case,  $S$  must couple at tree-level to the messenger fields in the superpotential (just as in the model of the previous section). Now, the issue is to prevent  $F_S$  from getting a VEV at the same loop order. Our main remark is that, since  $F_S$  contributes to the vacuum energy, it may be small just due to energetic reasons. The sufficient condition for this to happen is that  $F_S$  is a function of some field(s) which, at one-loop order, has no potential apart from  $F_S^* F_S$ . Thus the field(s) can slide and compensate any source for  $F_S$ . In our model such a field was  $N$ . In this case  $F_S$  and thus  $B$  can only appear as higher loop-effects and  $B\mu$  is acceptably small. Thus, the compensator field  $N$  must have no superpotential in the limit  $S = 0$  and we are lead to the following general structure for the superpotential (again we omit the  $\mathcal{O}(1)$  Yukawa couplings)

$$W = S(H\bar{H} + f(N) + \bar{Q}Q) + X\bar{Q}Q \quad (15)$$

In the model of the previous section we had  $f(N) = \lambda_2 N^2/2 - \lambda \Lambda^2$ , and the scale  $\Lambda$  came from the condensate of strongly coupled messengers  $\bar{Q}Q$ . Now we will study what is the possible general form of the  $f$ -function. For the time being we will be assuming a single fundamental scale  $F_X \sim X^2$ , although this is not in any respect necessary for our mechanism to work as it was clear from our model in which  $X^2 \gg F_X$ . The key idea is most transparent in units  $\langle F_X \rangle = 1$ . In these units, as we will now show, the function  $f$  must satisfy the following necessary conditions in the minimum to zeroth-loop order:

$$f = 0, \quad f' \sim 1, \quad f'' \sim 1 \quad (16)$$

where prime denotes derivative with respect to the compensator field  $N$ . Now since  $S$  is the only field that gets a correction to the potential from one-loop one-particle-irreducible diagrams, the effective potential to this order can be written as (both Higgs doublets and messengers are put to zero as it should be)

$$V = |f|^2 + |S|^2 |f'|^2 + V_1 \quad (17)$$

where  $V_1$  is a one-loop effective potential for  $S$ . Since the  $S$ -VEV in any case will be stabilized by the second term, which gives a curvature  $\sim 1$ , it is sufficient to keep only a

tadpole part linear in  $S$  in  $V_1$ , which is of order  $\epsilon S$ . In what follows  $\epsilon$  denotes a one-loop suppression factor.

Now, we minimize with respect to  $S$  and  $N$  remembering that  $B\mu = F_s = f$ , up to factors of order one, and we get the following equations

$$f^{*'} B\mu + f^{*''} f' |S|^2 = 0, \quad |f'|^2 s + \epsilon = 0 \quad (18)$$

since the derivatives of  $f$  are  $\sim 1$  in the zeroth order, the same should hold at one-loop, and, thus, we immediately get that  $\mu \sim S \sim \epsilon$  and  $B\mu \sim \epsilon^2$ .

The simplest explicit model with dynamical relaxation [18] is based on  $f = N^2 - \Lambda^2$  and we have shown in the previous section the scale  $\Lambda$  can be indeed generated dynamically by the condensate of the strongly coupled messenger mesons. What if the messengers are not transforming under a strong gauge group? It is interesting that in this case, the scale  $\Lambda$  in  $f$  can be induced through the kinetic mixing of  $X$  and  $S$  superfields in the Kähler potential. This is the analog of the term  $Z_{YS}$  in eq. 4. Such a mixed term is not forbidden by any symmetry and even if not present at the tree-level will be induced through the loops with  $Q$ -particles. Assuming  $Z_{XS}$  to be zero at the Planck scale  $M_P$ , the resulting mixed term in the Kähler potential has the form

$$\Delta K \sim \frac{n5\lambda\rho}{16\pi^2} \ln\left(\frac{M_P^2}{XX^*}\right) SX^* \quad (19)$$

where  $n$  is the number of messengers coupled to  $S$ . After substituting  $F_X \neq 0$  and solving the equation of motion of  $F_S$ , the effect of this term is just a shift in  $F_S$

$$F_S \rightarrow \lambda_2 N^2/2 + F_X \frac{n5\lambda\rho}{16\pi^2} \ln\left(\frac{M_P^2}{XX^*}\right) \quad (20)$$

The resulting  $\mu \sim \sqrt{B\mu}$  term that is generated through the vacuum relaxation mechanism is suppressed by  $1/\ln(M_P/X)$  factor with respect to the original supersymmetry-breaking scale. For instance the ratio between the squark mass and  $\mu$  reads

$$\frac{\mu}{m_{\tilde{Q}}} = \sqrt{\frac{3}{8n}} \frac{\lambda_1 \pi}{\lambda_2 \alpha_3 \ln(M_P/X)} \quad (21)$$

so that there seems to be a range of  $n$  and  $X$  where this ratio is close to 1.

Finally we will mention one more logical possibility which can dynamically introduce a scale of order  $F_X$  in the function  $f$  and satisfy conditions (16) necessary for the dynamical

relaxation mechanism. This is to have a SUSY breaking sector communicating with the compensators through the  $D$ -term of some gauge  $U(1)$ . Assume that the supersymmetry breaking VEVs develop along a  $D$ -non-flat direction of some  $U(1)$  gauge factor. This induces a non-zero expectation value of the  $D$ -term  $D \sim g$  (in  $F_X$  units) where  $g$  is a  $U(1)$  gauge-coupling constant. Let us introduce a pair of compensator superfields  $N_-$  and  $N_+$  with charges  $-1$  and  $+1$  under the above  $U(1)$ . The function  $f$  can then be chosen as  $f = N_- N_+$ . In the zeroth-loop order the  $N$ -fields have an  $F$ -flat potential and one of them, (say  $N_-$ ) can slide and compensate the  $D$ -term picking up a VEV of order one. The conditions (16) are then automatically satisfied. The  $\mu$  term is induced as a one-loop effect through the usual tadpole term, whereas  $B\mu$  is zero in this order and is only induced by higher loop corrections.

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## REFERENCES

- [1] M. Dine, W. Fischler, and M. Srednicki, Nucl. Phys. **B189** (1981) 575;  
S. Dimopoulos and S. Raby, Nucl. Phys. **B192** (1981) 353;  
M. Dine and W. Fischler, Phys. Lett. **B110** (1982) 227;  
M. Dine and M. Srednicki, Nucl. Phys. **B202** (1982) 238;  
M. Dine and W. Fischler, Nucl. Phys. **B204** (1982) 346;  
L. Alvarez-Gaumé, M. Claudson, and M. Wise, Nucl. Phys. **B207** (1982) 96;  
C.R. Nappi and B.A. Ovrut, Phys. Lett. **B113** (1982) 175;  
S. Dimopoulos and S. Raby, Nucl. Phys. **B219** (1983) 479.
- [2] S. Dimopoulos and D. Sutter, Nucl. Phys. **B452** (1995) 496.
- [3] M. Dine and A. Nelson, Phys. Rev. **D48** (1993) 1277; M. Dine, A. Nelson and Y. Shirman, Phys. Rev. **D51** (1995) 1362; M. Dine, A. Nelson, Y. Nir and Y. Shirman, Phys. Rev. **D53** (1996) 2658.
- [4] T. Hotta, K.-I. Izawa and T. Yanagida, Phys. Rev. **D55** (1997) 415.
- [5] I. Dasgupta, B.A. Dobrescu and L. Randall, Nucl. Phys. **B483** (1997) 95.
- [6] E. Poppitz and S. Trivedi, Phys. Rev. **D55** (1997) 5508.
- [7] N. Arkadi-Hamed, J. March-Russell, H. Murayama, preprint LBL-39865, hep-ph/9701286.
- [8] Y. Shadmi, preprint Fermilab-Pub-97/060-T, hep-ph/9703312.
- [9] H. Murayama, hep-ph/9705271.
- [10] S. Dimopoulos, G.Dvali, R.Rattazzi and G. Giudice, hep-ph/9705307.
- [11] M.A. Luty, hep-ph/9706554
- [12] S.L. Dubovskii, D.S. Gorbunov, S.V. Troitskii, hep-ph/9707357
- [13] C. Csaki, L. Randall, W. Skiba, hep-ph/9707386.
- [14] K. Intriligator and S. Thomas, Nucl. Phys. **B473** (1996) 121; K.-I. Izawa and

- T. Yanagida, *Prog. Theor. Phys.* **95** (1996) 829.
- [15] N. Seiberg, *Phys. Rev.* **D49** (1994) 6857. For a review see, e.g. K. Intriligator and N. Seiberg, hep-th/9509066.
  - [16] G.Dvali and A.Pomarol, *Phys. Rev. Lett.* **77** (1996) 3728; P.Binetruy and E. Dudas, *Phys. Lett.* **B389** (1996) 503.
  - [17] G.Giudice and R. Rattazzi, hep-ph/9706540.
  - [18] G. Dvali, G.F. Giudice, and A. Pomarol, *Nucl. Phys.* **B478** (1996) 31.
  - [19] E. Witten, *Phys. Lett.* **B105** (1981) 267.
  - [20] K. Inoue, A. Kakuto and T. Takano, *Prog. Theor. Phys.* **75** (1986) 664; A. Anselm and A. Johansen, *Phys. Lett.* **B200** (1988) 331; A. Anselm, *Sov. Phys. JETP* **67** (1988) 663; Z. Berezhiani and G. Dvali, *Sov. Phys. Lebedev Institute Reports* **5** (1989) 55; R. Barbieri, G. Dvali and M. Moretti, *Phys. Lett.* **B312** (1993) 137; Z. Berezhiani, C. Csaki and L. Randall, *Nucl. Phys.* **B444** (1995) 61; G. Dvali and S. Pokorski, *Phys. Rev. Lett.* **78** (1997) 807.
  - [21] P. Ciafaloni and A. Pomarol, *Phys. Lett.* **B404** (1997) 83.
  - [22] G. Veneziano and S. Yankielowicz, *Phys. Lett.* **B113** (1982) 321; T. R. Taylor, G. Veneziano and S. Yankielowicz, *Nucl. Phys.* **B218** (1983) 493.
  - [23] V. Novikov, M. Shifman, A. Vainshtein and V. Zakharov, *Nucl. Phys.* **B229** (1983) 381.
  - [24] M.A. Shifman, A.I. Vainstein, *Nucl. Phys.* **B277** (1986) 456; *Nucl. Phys.* **B359** (1991) 571.